

# Exclusive $B \rightarrow VV$ Decays and CP Violation in the General two-Higgs-doublet Model

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## Abstract

Using the general factorization approach, we present a detailed investigation for the branching ratios, CP asymmetries and longitudinal polarization fractions in all charmless hadronic  $B \rightarrow VV$  decays (except for the pure annihilation processes) within the most general two-Higgs-doublet model with spontaneous CP violation. It is seen that such a new physics model only has very small contributions to the branching ratios and longitudinal polarization fractions. However, as the model has rich CP-violating sources, it can lead to significant effects on the CP asymmetries, especially on those of penguin-dominated decay modes, which provides good signals for probing new physics beyond the SM in the future B-physics experiments.

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## I. INTRODUCTION

During the recent years, tremendous progress in B physics has been made through the fruitful interplay between theory and experiment. The precise measurements of the  $B$ -meson decays can provide an insight into very high energy scales via the indirect loop effects of new physics beyond the standard model (SM), which makes the study of exclusive non-leptonic  $B$ -meson decays of great interest.

In the SM, the phenomenon of CP violation can be accommodated in an efficient way through a complex phase entering the quark-mixing matrix, which governs the strength of the charged-current interactions of the quarks. This Kobayashi-Maskawa (KM) [1] mechanism of CP violation is the subject of detailed investigation in these few decades. However, its origin remains unknown as it is put into the standard model through the complex Yukawa couplings. Moreover, the baryon asymmetry of the universe requires new sources of CP violation. Many possible extensions of the SM in the Higgs sector have been proposed [2], and it was suggested that CP symmetry may break down spontaneously [3]. A consistent and simple model, which provides a spontaneous CP violation mechanism, has been constructed completely in a general two-Higgs-doublet model (2HDM) [4, 5] without imposing the *ad hoc* discrete symmetry, which is now commonly called as type III 2HDM. The type III 2HDM, which allows flavor-changing neutral currents (FCNCs) at tree level but suppressed by approximate  $U(1)$  flavor symmetry, has attracted much more interests. It is known that FCNCs are suppressed in low-energy experiments, especially for the lighter two generation quarks. Thus, the type III 2HDM can be parameterized in a way to satisfy the current experimental constraints. On the other hand, constraints on the general 2HDM from the neutral mesons mixing ( $K^0 - \bar{K}^0$ ,  $D^0 - \bar{D}^0$ , and  $B^0 - \bar{B}^0$ ) [6, 7] and from the radiative decays of bottom quark [8] have also been studied in detail.

In recent years, there are many works about the  $B$ -meson decays within the two-Higgs-doublet model. In Refs. [9, 10], the authors have studied the  $B \rightarrow PP, PV$  decays (with  $P$  and  $V$  denoting the pseudoscalar and vector mesons, respectively) within the type III 2HDM. Since through the measurements of magnitudes and phases of various helicity amplitudes, the charmless hadronic  $B \rightarrow VV$  decay modes can reveal more dynamics of exclusive  $B$  decays than  $B \rightarrow PP$  and  $B \rightarrow PV$  decays, in the present work we are going to make a detailed study for  $B \rightarrow VV$  decays within the type III 2HDM by emphasizing on the new physics contributions. It will be seen that this specific new physics has remarkable effects on CP asymmetries, especially on the parameter  $S_f$  for the penguin-dominated decay modes. On the other hand, the new physics is found to have very small contributions to the branching ratios and the transverse polarizations. Furthermore, the polarization anomaly observed in  $B \rightarrow \rho K^*$  and  $B \rightarrow \phi K^*$  modes can not be improved in our current considered parameter spaces.

The paper is organized as follows: In section II, we first describe the theoretical framework, including a brief introduction for the two-Higgs-doublet model with spontaneous CP violation, the effective Hamiltonian, as well as the decay amplitudes and CP violation formulas, which are the basic tools to estimate the branching ratios and CP asymmetries of  $B$ -meson decays. In section III, we list the Wilson coefficients and the other relevant input parameters. Our numerical predictions for the branching ratios, CP asymmetries and longitudinal polarization fractions are presented in Section IV. Our conclusions are presented in the last section.

## II. THEORETICAL FRAMEWORK

### A. Outline of the Two-Higgs-doublet Model

Motivated solely from the origin of CP violation, a general two-Higgs-doublet model with spontaneous CP violation (type III 2HDM) has been shown to provide one of the simplest and attractive models in understanding the origin and mechanism of CP violation at the weak scale. In such a model, there exists more physical neutral and charged Higgs bosons and rich CP violating sources from a single CP phase of the vacuum. These new sources of CP violation can lead to some significant phenomenological effects, which are promising to be tested by the future  $B$  factory and the LHCb experiments. In this paper, we shall focus on the phenomenological applications of the type III 2HDM on the two-body charmless hadronic  $B \rightarrow VV$  decays.

The two complex Higgs doublets in the general 2HDM are generally expressed as [4, 5, 12, 13]

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \phi_1^0 \end{pmatrix}, \quad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \phi_2^0 \end{pmatrix}. \quad (1)$$

The corresponding Yukawa Lagrangian is given as

$$\mathcal{L}_Y = \eta_{ija} \bar{\psi}_{i,L} \tilde{\Phi}_a U_{j,R} + \xi_{ija} \bar{\psi}_{i,L} \Phi_a D_{j,R} + h.c., \quad (2)$$

where the parameters  $\eta_{ija}$  and  $\xi_{ija}$  are real, so that the lagrangian is CP invariant. After the symmetry is spontaneously broken down

$$\langle \phi_1^0 \rangle = v_1 e^{i\alpha_1}, \quad \langle \phi_2^0 \rangle = v_2 e^{i\alpha_2}, \quad (3)$$

and the Goldstone particles have been eaten, the physical Higgs bosons are

$$H_1 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + \phi_1^0 \end{pmatrix}, \quad H_2 = \frac{1}{\sqrt{2}} \begin{pmatrix} H^+ \\ \phi_2^0 + i\phi_3^0 \end{pmatrix}. \quad (4)$$

where  $H^\pm$  are the charged scalar mass eigenstates,  $(\phi_1^0, \phi_2^0, \phi_3^0)$  are generally not the mass eigenstates but can be expressed as linear combinations of the mass eigenstates  $(H, h, A)$ .

Then the Yukawa part of the Lagrangian for physical particles can be written as

$$\mathcal{L}_Y = \eta_{ij}^U \bar{\psi}_{i,L} \tilde{H}_1 U_{j,R} + \eta_{ij}^D \bar{\psi}_{i,L} H_1 D_{j,R} + \xi_{ij}^U \bar{\psi}_{i,L} \tilde{H}_2 U_{j,R} + \xi_{ij}^D \bar{\psi}_{i,L} H_2 D_{j,R} + h.c., \quad (5)$$

where

$$\begin{aligned} \eta_{ij}^U &= \eta_{ij1} \cos \beta + \eta_{ij2} e^{-\delta} \sin \beta, \\ \xi_{ij}^U &= -\eta_{ij1} e^{-\delta} \sin \beta + \eta_{ij2} \cos \beta, \\ \eta_{ij}^D &= \xi_{ij1} \cos \beta + \xi_{ij2} e^{-\delta} \sin \beta, \\ \xi_{ij}^D &= -\xi_{ij1} e^{-\delta} \sin \beta + \xi_{ij2} \cos \beta, \end{aligned} \quad (6)$$

and these couplings  $\eta^U, \eta^D, \xi^U, \xi^D$  are generally complex, which means CP violation. According to the CKM mechanism, after diagonalizing the fermion terms' couplings  $\eta^U$  and

$\eta^D$ , the other couplings become

$$\begin{aligned}\mathcal{L}_Y = & \bar{U}_i \frac{m^U}{v} U_R(v + \phi_1^0) + \bar{D}_L \frac{m^D}{v} D_R(v + \phi_1^0) \\ & + \bar{U}_L \tilde{\xi}^U U_R(\phi_2^0 + i\phi_3^0) + \bar{D}_L \hat{\xi}^U U_R H^- \\ & + \bar{U}_L \hat{\xi}^D D_R H^+ + \bar{D}_L \tilde{\xi}^D D_R(\phi_2^0 + i\phi_3^0) + h.c.,\end{aligned}\tag{7}$$

with

$$\begin{aligned}\xi^{U,D} &= (V_L^{U,D})^{-1} \xi^{U,D} V_R^{U,D}, \\ \hat{\xi}^U &= \tilde{\xi}^U V_{\text{CKM}}, \\ \hat{\xi}^D &= V_{\text{CKM}} \tilde{\xi}^D.\end{aligned}\tag{8}$$

The Yukawa couplings may be parameterized as following

$$\tilde{\xi}_{ij} = \lambda_{ij} \frac{\sqrt{m_i m_j}}{v}.\tag{9}$$

with  $v$  the vacuum expectation value  $v = 246$  GeV.

## B. Effective Hamiltonian and decay amplitudes of $B \rightarrow VV$ decays

Using the operator product expansion and the renormalization group equation, the low energy effective Hamiltonian for charmless hadronic  $B$ -meson decays with  $\Delta B = 1$  can be written as

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}} \sum_{p=u,c} V_{pb} V_{pq}^* \left( C_1 Q_1^p + C_2 Q_2^p + \sum_{i=3,\dots,16} [C_i Q_i + C'_i Q'_i] \right) + h.c.,\tag{10}$$

where  $C_i(\mu)$  ( $i = 1, \dots, 16$ ) are the Wilson coefficients that can be calculated by perturbative theory, and  $Q_i$  are the quark and gluon effective operators, with  $Q_{1-10}$  and  $Q_{11-16}$  coming from the SM and from the type III 2HDM, respectively. Their explicit forms are

defined as follows (taking  $b \rightarrow sq\bar{q}$  transition as an example) [17]

$$\begin{aligned}
Q_1 &= (\bar{s}u)_{V-A}(\bar{u}b)_{V-A}, \\
Q_2 &= (\bar{s}_i u_j)_{V-A}(\bar{u}_j b_i)_{V-A}, \\
Q_{3(5)} &= (\bar{s}b)_{V-A} \sum_q (\bar{q}q)_{V-(+)A}, \\
Q_{4(6)} &= (\bar{s}_i b_j)_{V-A} \sum_q (\bar{q}_j q_i)_{V-(+)A}, \\
Q_{7(9)} &= \frac{3}{2}(\bar{s}b)_{V-A} \sum_q e_q (\bar{q}q)_{V+(-)A}, \\
Q_{8(10)} &= \frac{3}{2}(\bar{s}_i b_j)_{V-A} \sum_q e_q (\bar{q}_j q_i)_{V+(-)A}, \\
Q_{11(13)} &= (\bar{s}b)_{S+P} \sum_q \frac{m_q \lambda_{qq}^* (\lambda_{qq})}{m_b} (\bar{q}q)_{S-(+)P}, \\
Q_{12(14)} &= (\bar{s}_i b_j)_{S+P} \sum_q \frac{m_q \lambda_{qq}^* (\lambda_{qq})}{m_b} (\bar{q}_j q_i)_{S-(+)P}, \\
Q_{15} &= \bar{s} \sigma^{\mu\nu} (1 + \gamma_5) b \sum_q \frac{m_q \lambda_{qq}}{m_b} \bar{q} \sigma_{\mu\nu} (1 + \gamma_5) q, \\
Q_{16} &= \bar{s}_i \sigma^{\mu\nu} (1 + \gamma_5) b_j \sum_q \frac{m_q \lambda_{qq}}{m_b} \bar{q}_j \sigma_{\mu\nu} (1 + \gamma_5) q_i,
\end{aligned} \tag{11}$$

where  $(\bar{q}_1 q_2)_{V\pm A} = \bar{q}_1 \gamma_\mu (1 \pm \gamma_5) q_2$  and  $(\bar{q}_1 q_2)_{S\pm P} = \bar{q}_1 (1 \pm \gamma_5) q_2$ , with  $qu, d, s, c, b$ , and  $e_q$  is the electric charge number of  $q$  quark. The operators  $Q'_i$  in Eq. (11) are obtained from  $Q_i$  via exchanging  $L \leftrightarrow R$ , and we shall neglect their effects in our calculations for they are suppressed by a factor  $m_s/m_b$  in model III 2HDM. The Wilson coefficients  $C_i$  ( $i = 1, \dots, 10$ ) have been calculated at leading order (LO) [14, 15] and at next-to-leading order (NLO) [16] in the SM and also at LO in 2HDM [12], while  $C_i$  ( $i = 11, \dots, 16$ ) at LO can be found in Refs. [17, 18].

Having defined the effective Hamiltonian  $H_{eff}$  in terms of the four-quark operators  $Q_i$ , we can then proceed to calculate the hadronic matrix elements with the generalized factorization assumption [19, 20, 21, 22] based on the naive factorization approach.

For two-body charmless hadronic  $B \rightarrow VV$  decays, the decay amplitude of the local four fermion operators is defined as

$$A_h \equiv \frac{G_F}{\sqrt{2}} \langle V_1(h_1) V_2(h_2) | (\bar{q}_2 q_3)_{V\pm A} (\bar{b} q_1)_{V-A} | B \rangle, \tag{12}$$

where  $h_1$  and  $h_2$  are the helicities of the final-state vector mesons  $V_1$  and  $V_2$  with four-momentum  $p_1$  and  $p_2$ , respectively. Since the  $B$  meson has spin zero, in the rest frame of  $B$ -meson system, the two vector mesons have the same helicity due to helicity conservation. Therefore three polarization states are possible in  $B \rightarrow VV$  decays with one longitudinal ( $L$ ) and two transverse, corresponding to helicities  $h = 0$  and  $h = \pm$  (here

$h_1 = h_2 = h$ ), respectively. We define the three helicity amplitudes as follows

$$\begin{aligned} A_0 &= A(B \rightarrow V_1(p_1, \epsilon_1^0) V_2(p_2, \epsilon_2^0)), \\ A_{\pm} &= A(B \rightarrow V_1(p_1, \epsilon_1^{\pm}) V_2(p_2, \epsilon_2^{\pm})). \end{aligned} \quad (13)$$

We choose the momentum  $\vec{p}_2$  to be directed in the positive  $z$ -direction in the  $B$ -meson rest frame, and the polarization four-vectors of the light vector mesons such that in a frame where both light mesons have large momentum along the  $z$ -axis. They are given by

$$\begin{aligned} \epsilon_1^{\pm\mu} &= \epsilon_2^{\mp\mu} = (0, \pm 1, i, 0)/\sqrt{2}, \\ \epsilon_{1,2}^{0\mu} &= p_{1,2}^{\mu}/m_{1,2}, \end{aligned} \quad (14)$$

where  $m_1$  and  $m_2$  are the masses of  $V_1$  and  $V_2$  mesons, respectively. Using the definitions for decay constants and form factors [23], the tree-level hadronic matrix elements of the effective operators  $Q_i$  can be decomposed as the following two amplitudes

$$A_h = \mathcal{V}_h + \mathcal{T}_h, \quad (15)$$

with

$$\begin{aligned} \mathcal{V}_h &\equiv \langle V_1(p_1, \epsilon_1^h) | V - A | B \rangle \langle V_2(p_2, \epsilon_2^h) | V - A | 0 \rangle, \\ \mathcal{T}_h &\equiv \langle V_1(p_1, \epsilon_1^h) | \sigma^{\mu\nu} (1 + \gamma^5) | B \rangle \langle V_2(p_2, \epsilon_2^h) | \sigma_{\mu\nu} (1 + \gamma^5) | 0 \rangle. \end{aligned} \quad (16)$$

Here, for simplicity, we have omitted the quark spinors in the corresponding current operators in the above definitions. The three polarization amplitudes for  $\mathcal{V}^h$  and  $\mathcal{T}^h$  can be further written as

$$\begin{aligned} \mathcal{V}_0 &= i f_{V_2} (m_B^2 - m_1^2 - m_2^2) A_0^{V_1}, \\ \mathcal{V}_{\pm} &= i f_{V_2} m_2 \left[ A_1^{V_1} (m_1 + m_B) \mp V^{V_1} \frac{2m_B |p_c|}{m_B + m_1} \right], \\ \mathcal{T}_0 &= 0, \\ \mathcal{T}_{\pm} &= 2i f_{V_2}^{\perp} \left[ 2T_1^{V_1} m_B |p_c| \mp T_2^{V_1} (m_B^2 - m_1^2) \right]. \end{aligned} \quad (17)$$

From the amplitude given by Eq. (15), the branching ratio for  $B \rightarrow VV$  decays then reads

$$Br(B \rightarrow VV) = \frac{\tau_B |p_c|}{8\pi m_B^2} (|A_0|^2 + |A_+|^2 + |A_-|^2), \quad (18)$$

where  $\tau_B$  is the lifetime of the  $B$  meson, and  $p_c$  is the center of mass momentum of either final-state meson with

$$|p_c| = \frac{\sqrt{[m_B^2 - (m_1 + m_2)^2][m_B^2 - (m_1 - m_2)^2]}}{2m_B}. \quad (19)$$

In order to compare the relative size of the three different helicity amplitudes, we can define the longitudinal polarization fraction as

$$f_L = \frac{|A_0|^2}{|A_0|^2 + |A_+|^2 + |A_-|^2}, \quad (20)$$

which measures the relative strength of the longitudinally polarization amplitude in a given decay mode.

### C. CP-violating asymmetries in $B \rightarrow VV$ decays

Since there are abundant CP violation sources in the two-Higgs-doublet model, it is also necessary and interesting for us to discuss CP asymmetries in  $B \rightarrow VV$  decays.

Firstly, for charged  $B^\pm$ -meson decays, there is only one simple type of CP violating asymmetry, which detects direct CP violation

$$\mathcal{A}_{CP} \equiv \frac{\Gamma(B^+ \rightarrow f^+) - \Gamma(B^- \rightarrow f^-)}{\Gamma(B^+ \rightarrow f^+) + \Gamma(B^- \rightarrow f^-)} \quad (21)$$

For neutral  $B$ -meson decays, there is another type of CP violation coming from the mixing between  $B_q^0 - \bar{B}_q^0$  (here  $q = d$  or  $s$ )

$$\begin{aligned} |B_q^0(t)\rangle &= g_+(t)|B_q^0\rangle + \frac{q}{p}g_-(t)|\bar{B}_q^0\rangle, \\ |\bar{B}_q^0(t)\rangle &= \frac{p}{q}g_-(t)|B_q^0\rangle + g_+|\bar{B}_q^0\rangle. \end{aligned} \quad (22)$$

In this case, there are in general four amplitudes which can be expressed as [24, 25, 26]

$$\begin{aligned} A_f &= \langle f | H_{eff} | B_q^0 \rangle, \quad \bar{A}_f = \langle f | H_{eff} | \bar{B}_q^0 \rangle, \\ \bar{A}_{\bar{f}} &= \langle \bar{f} | H_{eff} | \bar{B}_q^0 \rangle, \quad A_{\bar{f}} = \langle \bar{f} | H_{eff} | B_q^0 \rangle. \end{aligned} \quad (23)$$

For the  $B_d - \bar{B}_d$  and  $B_s - \bar{B}_s$  systems, the following approximations can be made

$$\text{both } B_d \text{ and } B_s \text{ systems : } \left| \frac{q}{p} \right| \sim 1; \quad \text{only } B_d \text{ system : } \Delta\Gamma \sim 0. \quad (24)$$

Using the decay amplitudes and the approximations listed in Eqs. (23) and (24), the time-dependent decay probabilities for  $B_d$  system can then be written as

$$\begin{aligned} \Gamma(B_d^0(t) \rightarrow f) &= \frac{|A_f|^2(1 + |\lambda_f|^2)}{2} e^{-\Gamma t} \{1 + C_f \cos(\Delta m t) - S_f \sin(\Delta m t)\}, \\ \Gamma(\bar{B}_d^0(t) \rightarrow f) &= \frac{|A_f|^2(1 + |\lambda_f|^2)}{2} e^{-\Gamma t} \{1 - C_f \cos(\Delta m t) + S_f \sin(\Delta m t)\}, \end{aligned} \quad (25)$$

while for  $B_s$  system, we have

$$\begin{aligned} \Gamma(B_s^0(t) \rightarrow f) &= \frac{|A_f|^2(1 + |\lambda_f|^2)}{2} e^{-\Gamma t} \left[ \cosh\left(\frac{\Delta\Gamma t}{2}\right) + D_f \sinh\left(\frac{\Delta\Gamma t}{2}\right) \right. \\ &\quad \left. + C_f \cos(\Delta m t) - S_f \sin(\Delta m t) \right], \\ \Gamma(\bar{B}_s^0(t) \rightarrow f) &= \frac{|A_f|^2(1 + |\lambda_f|^2)}{2} e^{-\Gamma t} \left[ \cosh\left(\frac{\Delta\Gamma t}{2}\right) + D_f \sinh\left(\frac{\Delta\Gamma t}{2}\right) \right. \\ &\quad \left. - C_f \cos(\Delta m t) + S_f \sin(\Delta m t) \right], \end{aligned} \quad (26)$$

where  $\Gamma$  is the average decay width,  $\Delta\Gamma$  and  $\Delta m$  are the width and mass difference, respectively. The other quantities are defined as

$$\begin{aligned} \lambda_f &\equiv \frac{q}{p} \frac{\bar{A}_f}{A_f}, & D_f &\equiv \frac{2\text{Re}(\lambda_f)}{1 + |\lambda_f|^2}, \\ C_f &\equiv \frac{1 - |\lambda_f|^2}{1 + |\lambda_f|^2}, & S_f &\equiv \frac{2\text{Im}(\lambda_f)}{1 + |\lambda_f|^2}. \end{aligned} \quad (27)$$

From Eqs.(25) and (26), we can get:

$$\begin{aligned}\mathcal{A}_{\mathcal{CP}}(B_d \rightarrow f) &= -C_f \cos \Delta mt + S_f \sin \Delta mt, \\ \mathcal{A}_{\mathcal{CP}}(B_s \rightarrow f) &= \frac{-C_f \cos \Delta mt + S_f \sin \Delta mt}{\cosh\left(\frac{\Delta \Gamma t}{2}\right) + D_f \sinh\left(\frac{\Delta \Gamma t}{2}\right)}.\end{aligned}\quad (28)$$

### III. INPUT PARAMETERS

The theoretical predictions in our calculations depend on many input parameters, such as the Wilson coefficients, the CKM matrix elements, the hadronic parameters, and so on. Here we present all the relevant input parameters as follows.

It has been shown from  $B_{d,s}^0 - \bar{B}_{d,s}^0$  mixings that the parameters  $|\lambda_{cc}|$  and  $|\lambda_{ss}|$  in Eq. (11) can reach to be around 100 [27], while their phases are not well constrained. In our present work we simply fix the phases to be  $\pi/4$ , and this choice will not cause any trouble in our numerical results. For the parameters  $\lambda_{tt}$  and  $\lambda_{bb}$ , the constraints come mainly from the experiments for  $B - \bar{B}$  mixing,  $\Gamma(b \rightarrow s\gamma)$ ,  $\Gamma(b \rightarrow c\tau\bar{\nu}_\tau)$ ,  $\rho_0$ ,  $R_b$ ,  $B \rightarrow PV$ , and the electric dipole moments (EDMS) of the electron and neutron [10, 11, 12, 17, 18]. Based on the above analyses, we choose the following three typical parameter spaces which are allowed by the present experiments and have been adopted for the  $B \rightarrow PV$  decays[10]

$$\begin{aligned}\text{Case A : } & |\lambda_{tt}| = 0.15; \quad |\lambda_{bb}| = 50, \\ \text{Case B : } & |\lambda_{tt}| = 0.3; \quad |\lambda_{bb}| = 30, \\ \text{Case C : } & |\lambda_{tt}| = 0.03; \quad |\lambda_{bb}| = 100,\end{aligned}$$

and  $\theta_{tt} + \theta_{bb} = \pi/2$ . For the Higgs masses and the Wilson coefficients of  $C_{1,\dots,10}$  corresponding to the SM, we use the results listed in the paper [10], while for the Wilson coefficients in the type III 2HDM, we redefine them as  $\tilde{C}_{11,\dots,16} = \frac{m_s \lambda_{ss}^{(*)}}{m_b} C_{11,\dots,16}$  in order to compare the contributions from those operators in SM, here the factor  $\frac{m_s \lambda_{ss}^{(*)}}{m_b}$  is associated with the operators in 2HDM, the numerical values for  $\tilde{C}_{11,\dots,16}$  are listed in Table I.

TABLE I: The Wilson coefficients  $\tilde{C}_{11,\dots,16} = \frac{m_s \lambda_{ss}^{(*)}}{m_b} C_{11,\dots,16}$  in  $b \rightarrow s$  transition at  $\mu = m_b = 4.2$  GeV in 2HDM.

	Case A	Case B	Case C
$\tilde{C}_{11}$	$-0.0085 + 0.012i$	$-0.0085 + 0.018i$	$-0.010 + 0.012i$
$\tilde{C}_{12}$	0	0	0
$\tilde{C}_{13}$	$-0.0030 - 0.0049i$	$-0.0052 - 0.0069i$	$-0.0029 - 0.0052i$
$\tilde{C}_{14}$	$-0.000060 - 0.00010i$	$-0.00011 - 0.00014i$	$-0.000059 - 0.00010i$
$\tilde{C}_{15}$	$0.000033 + 0.000055i$	$0.000058 + 0.000078i$	$0.000032 + 0.000059i$
$\tilde{C}_{16}$	$-0.00010 - 0.00017i$	$-0.00018 - 0.00024i$	$-0.0001 - 0.00018i$

As for the CKM matrix elements, we shall use the Wolfenstein parametrization [28] with the values [26]:  $A = 0.8533 \pm 0.0512$ ,  $\lambda = 0.2200 \pm 0.0026$ ,  $\bar{\rho} = 0.20 \pm 0.09$ , and  $\bar{\eta} = 0.33 \pm 0.05$ .



For the hadronic parameters, the decay constants, and the form factors, we list them in Tables. II and III, respectively.

TABLE II: The hadronic input parameters [26] and the decay constants taken from the QCD sum rules [29] and Lattice theory [30].

$\tau_{B^\pm}$	$\tau_{B_d}$	$\tau_{B_s}$	$M_{B_d}$	$M_{B_s}$	$m_b$
1.638ps	1.528ps	1.472ps	5.28GeV	5.37GeV	4.2GeV
$m_t$	$m_u$	$m_d$	$m_c$	$m_s$	$m_{\rho^0}$
174GeV	3.2MeV	6.4MeV	1.1GeV	0.105GeV	0.77GeV
$m_{\rho^\pm}$	$m_\omega$	$m_\phi$	$m_{K^{*\pm}}$	$m_{K^{*0}}$	$\Lambda_{QCD}$
0.77GeV	0.782GeV	1.02GeV	0.892GeV	0.896GeV	225MeV
$f_\rho$	$f_\omega$	$f_{K^*}$	$f_\phi$	$f_\rho^T$	$f_\omega^T$
0.205GeV	0.195GeV	0.217GeV	0.231GeV	0.147GeV	0.133GeV
$f_{K^*}^T$	$f_\phi^T$				
0.156GeV	0.183GeV				

TABLE III: The relevant  $B \rightarrow V$  transition form factors at  $q^2 = 0$  taken from the light-cone sum rules (LCSR) [31, 32].

decay channel	$V$	$A_0$	$A_1$	$A_2$	$T_1$	$T_3$
$B \rightarrow \rho$	0.323	0.303	0.242	0.221	0.267	0.176
$B \rightarrow \omega$	0.293	0.281	0.219	0.198	0.242	0.155
$B \rightarrow K^*$	0.411	0.374	0.292	0.259	0.333	0.202
$B_s \rightarrow \bar{K}^*$	0.311	0.363	0.233	0.181	0.26	0.136
$B_s \rightarrow \phi$	0.434	0.474	0.311	0.234	0.349	0.175

#### IV. NUMERICAL RESULTS AND DISCUSSIONS

In this section, we shall classify the 28 channels of  $B^+$ ,  $B^0$  and  $B_s$  decays into two light vector mesons according to the reliability of the calculation for various observables, which is motivated by the dominated contributing operators. We shall give our predictions for the branching ratios, the CP asymmetries, and the longitudinal polarization fractions both in the SM and in the 2HDM. Comparisons with the current experiment data, if possible, are also made.

Before moving to the detailed discussions, some general observations of new physics effects on  $B \rightarrow VV$  decays should be made. As can be seen from Eqs. (9) and (11), the contributions of new physics operators  $O_{11,\dots,16}$  are always proportional to the factor  $m_q/v$ . Thus, they are severely suppressed for the first generation quarks. In this case, for  $B \rightarrow \rho K^*, \omega K^*, \rho\rho, \omega\rho, \omega\omega$  and  $B_s \rightarrow \rho K^*, \omega K^*, K^* K^*, \rho\phi, \omega\phi$  decay channels, we can

safely ignore the contributions from those new operators. Note that the new physics still has effects on the Wilson Coefficients  $C_{1-10}$ . On the other hand, for  $B \rightarrow \phi K^*, \phi \rho, \phi \omega$  and  $B_s \rightarrow \phi K^*, \phi \phi$  decay channels, since these are all induced by  $b \rightarrow q s \bar{s}$  ( $q = d, s$ ) transitions, we could not ignore the new operators' contributions any more in this case. In the general factorization approach, it is impossible to produce a vector meson via the scalar and/or pseudoscalar currents from the vacuum state, and hence the new operators  $Q_{11}$  and  $Q_{13}$  have no contributions to  $B \rightarrow VV$  decays. Moreover, from the results listed in Table I, it can be seen that all the contributing new operators  $Q_{12,14,15,16}$  have only very small (even zero) Wilson coefficients. It is therefore expected that the new physics will have very small effects on the branching ratios and transverse amplitudes (hence on the transverse polarization fractions) of  $B \rightarrow VV$  decays.

### A. CP-averaged branching ratios and direct CP violation

According to different decay modes, we shall give our predictions for the branching ratios and direct CP violations one by one.

(i), color-allowed tree-dominated decays. Our predictions for the CP-averaged branching ratios and the direct CP asymmetries are presented in Table IV. From the numerical results, we can see that the branching ratios are all at  $10^{-5}$  order, and the direct CP asymmetries are all very small since the penguin amplitude contributions are much smaller than the ones from the tree diagrams. Most predictions within the SM are consistent with the current experiment data, and the new physics has very small effects on these types of decays.

TABLE IV: The CP-averaged branching ratios (in unit of  $10^{-6}$ ) (first line) and the direct CP violations (second line) for the color-allowed tree-dominant processes both in the SM and in the type III 2HDM. Case A-C stand for the three different parameter spaces listed in Section III.

Decay modes	Case A	Case B	Case C	SM	Exp.
$B^+ \rightarrow \rho^+ \rho^0$	14.59	14.59	14.59	15.53	$18.2 \pm 3.0$
	-0.004	-0.004	-0.004	-0.002	$-0.08 \pm 0.13$
$B^0 \rightarrow \rho^+ \rho^-$	26.33	25.93	26.73	27.49	$24.2^{+3.1}_{-3.2}$
	-0.043	-0.043	-0.042	-0.035	
$B^+ \rightarrow \rho^+ \omega$	12.66	12.47	12.85	13.97	$10.6^{+2.6}_{-2.3}$
	-0.042	-0.043	-0.042	-0.034	$0.04 \pm 0.18$
$B_s \rightarrow \rho^+ K^{*-}$	36.88	36.32	37.44	38.50	
	-0.043	-0.043	-0.042	-0.035	

(ii), color-suppressed tree-dominated decays. The numerical results are given in Table V, it is interesting to note that the branching ratios will generally become smaller after including the new physics contributions except for the  $B \rightarrow \rho^0 \rho^0$  mode. Furthermore, there are big direct CP violations in these decay processes except for the  $B^+ \rightarrow \rho^0 \omega$  mode, and the new physics has more effects on the direct CP asymmetries than on the

branching ratios through the Wilson coefficient functions, although there are no new operator contributions to the hadronic matrix elements in this type decays within our approximations. Compared to Case A and Case C, Case B has the biggest corrections to the CP asymmetries of the SM.

TABLE V: The same as Table IV but for color-suppressed tree-dominant processes.

Decay modes	Case A	Case B	Case C	SM	Exp.
$B^0 \rightarrow \rho^0 \rho^0$	0.0814	0.0897	0.0754	0.065	$0.86 \pm 0.28$
	0.176	0.218	0.119	0.153	
$B^+ \rightarrow \omega \omega$	0.112	0.110	0.115	0.160	$< 4.0$
	-0.117	-0.088	-0.144	-0.207	
$B_s \rightarrow \rho^0 \bar{K}^{*0}$	0.081	0.090	0.073	0.092	$< 7.67 \times 10^{-4}$
	0.176	0.218	0.119	0.153	
$B_s \rightarrow \omega \bar{K}^{*0}$	0.183	0.180	0.187	0.262	
	-0.167	-0.088	-0.144	-0.207	
$B^+ \rightarrow \rho^0 \omega$	0.024	0.024	0.024	0.076	$< 1.5$
	-0.063	-0.063	-0.063	-0.035	

(iii), penguin-dominated decays. We may divide such decays into two types:  $\Delta S = 1$  and  $\Delta D = 1$  decay modes. They are corresponding to the upper and the lower parts in Table VI, respectively. From the numerical results, we can see that all the eleven  $\Delta S = 1$  decay modes have branching ratios up to  $10^{-6}$  or even to  $10^{-5}$  order, since they involve the relative large CKM matrix elements  $V_{ts}^*$ , while the  $\Delta D = 1$  ones have much smaller branching ratios of order of  $10^{-7}$  due to the smaller CKM matrix elements  $V_{td}^*$ . For  $B \rightarrow \omega K^*$  and  $B_s \rightarrow \phi \phi$  decay modes, our predictions for the branching ratios with including the new operator contributions have similar results as the ones within the SM, which, however, are not quite consistent with the current experimental data; the numerical results for  $B \rightarrow \omega K^*$  modes are larger than the current experiment limit, and the prediction for  $B_s \rightarrow \phi \phi$  is about two times larger than the present data. For the other decay modes, our predictions for the branching ratios are in general agreement with the data. As for the direct CP asymmetries, there are big CP violations in some decay modes, and the new physics can lead to remarkable effects. Our predictions are consistent to the data in all these decay modes.

(iv), electroweak penguin or QCD flavor singlet dominated decays. As can be seen from Table VII, this type of decays are expected to have smaller branching ratios due to the large cancelations among the different Wilson coefficients. Although there are new operator contributions in  $B \rightarrow \rho \phi$  and  $\omega \phi$  decay modes, the predicted branching ratios are still small. The direct CP asymmetries for these decays are all small, and the new physics effects on these observables are not prominent. Due to the lack of accurate experimental data, we couldn't compare our predictions with the data yet.

(v), the pure annihilation decays. Only six decays belong to this class, namely  $B^0 \rightarrow K^{*+} K^{*-}$ ,  $B^0 \rightarrow \phi \phi$ ,  $B_s \rightarrow \rho^+ \rho^-$ ,  $B_s \rightarrow \rho^0 \rho^0$ ,  $B_s \rightarrow \rho^0 \omega$ , and  $B_s \rightarrow \omega \omega$ . Due to the lack of the information for the  $V_1 \rightarrow V_2$  transition form factor at large momentum transfers, we shall not consider them in details in this paper.

TABLE VI: The same as Table IV but for the penguin-dominated decay modes. The upper and the lower parts correspond to  $\Delta S = 1$  and  $\Delta D = 1$  processes, respectively.

Decay modes	Case A	Case B	Case C	SM	Exp.
$B^+ \rightarrow \rho^+ K^{*0}$	7.169	7.409	7.027	7.287	$9.2 \pm 1.5$
	0.084	0.117	0.049	0.018	$-0.01 \pm 0.16$
$B^+ \rightarrow \rho^0 K^{*+}$	5.853	6.229	5.526	5.575	$< 6.1$
	0.184	0.196	0.169	0.122	$0.20^{+0.32}_{-0.29}$
$B^0 \rightarrow \rho^0 K^{*0}$	6.396	6.513	6.324	6.245	$5.6 \pm 1.6$
	0.054	0.073	0.033	0.018	$0.09 \pm 0.19$
$B^0 \rightarrow \rho^- K^{*+}$	6.046	6.738	5.445	5.571	$< 12$
	0.295	0.301	0.283	0.199	
$B^0 \rightarrow \omega K^{*0}$	3.412	3.513	3.351	3.498	$< 2.7$
	0.078	0.107	0.048	0.024	
$B^+ \rightarrow \omega K^{*+}$	3.247	3.5697	2.965	3.123	$< 3.4$
	0.265	0.274	0.251	0.176	
$B^0 \rightarrow \phi K^{*0}$	9.276	9.704	9.221	9.318	$9.5 \pm 0.8$
	0.045	0.081	-0.002	0.020	$-0.01 \pm 0.06$
$B^+ \rightarrow \phi K^{*+}$	9.867	10.32	9.775	9.979	$10.0 \pm 1.1$
	0.039	0.074	-0.013	0.020	$-0.01 \pm 0.08$
$B_s \rightarrow \phi \phi$	28.99	30.34	28.64	28.85	$14^{+8}_{-7} \times 10^{-6}$
	0.054	0.089	0.006	0.020	
$B_s \rightarrow \bar{K}^{*0} K^{*0}$	9.303	9.614	9.118	9.456	$< 1.681 \times 10^{-3}$
	0.084	0.117	0.049	0.018	
$B_s \rightarrow K^{*+} K^{*-}$	8.404	9.366	7.569	7.744	
	0.295	0.302	0.283	0.199	
$B^0 \rightarrow \bar{K}^{*0} K^{*0}$	0.410	0.420	0.413	0.408	$0.49^{+0.17}_{-0.14}$
	-0.092	-0.061	-0.133	-0.145	
$B^+ \rightarrow K^{*+} K^{*0}$	0.439	0.450	0.443	0.437	$< 2.2$
	-0.092	-0.061	-0.133	-0.145	
$B_s \rightarrow \phi \bar{K}^{*0}$	0.517	0.532	0.521	0.526	$< 1.013 \times 10^{-3}$
	-0.094	-0.056	-0.145	-0.161	

## B. Time-dependent CP violating parameters $C_f$ , $S_f$ and $D_f$

Since there are abundant CP violating sources in type III 2HDM, it is expected that there are relatively large CP violations in 2HDM than in the SM. Using the relevant formulas given in section II, we can predict the time-dependent CP asymmetries in neutral  $B_d$  and  $B_s$  decays, with the numerical results given in Tables VIII and IX, respectively.

From these two tables, it is seen that, for  $B^0 \rightarrow \rho^+ \rho^-$ ,  $\rho^0 \phi$  and  $\omega \phi$  decay modes, the new physics has hardly any effects on the parameters  $C_f$  and  $S_f$ , even though there are new

TABLE VII: The same as Talbe IV but for the electroweak penguin or QCD flavor singlet dominated decays.

Decay modes	Case A	Case B	Case C	SM	Exp.
$B^+ \rightarrow \rho^+ \phi$	0.0054	0.0054	0.0054	0.0043	$< 16$
	-0.011	-0.011	-0.011	-0.014	
$B^0 \rightarrow \rho^0 \phi$	0.0025	0.0025	0.0025	0.0020	$< 13$
	-0.011	-0.011	-0.011	-0.014	
$B^0 \rightarrow \omega \phi$	0.0022	0.0022	0.0022	0.0017	$< 1.2$
	-0.011	-0.011	-0.011	-0.014	
$B_s \rightarrow \rho^0 \phi$	0.796	0.796	0.796	0.687	$< 6.17 \times 10^{-4}$
	0.0048	0.0048	0.0048	0.0039	
$B_s \rightarrow \phi \omega$	0.038	0.038	0.038	0.045	
	0.020	0.020	0.020	0.018	

operators contributions in  $B^0 \rightarrow \rho^0 \phi$  and  $\omega \phi$  decay modes. On the other hand, the new physics has remarkable effects on the other decay modes, especially on  $B^0 \rightarrow \omega \omega$  one (for this mode the new physics can even change the sign of the parameter  $S_f$ ). Furthermore, different parameter spaces also have remarkable effects on these CP violation parameters.

For  $B_s$  system, there are new operator contributions only in  $B_s \rightarrow \phi \phi$  mode. As is expected, the new physics has remarkable influence on the parameters  $C_f$ ,  $S_f$ , and  $D_f$ . For the other four decay modes, although there are no new operator contributions, the new physics still has big effects on the parameter  $S_f$ , but small effects on  $C_f$  and  $D_f$ .

### C. The polarization in $B \rightarrow \rho K^*$ and $\phi K^*$ decays

Motivated by the polarization anomaly observed by the BarBar [33], Belle [34] and CDF [35] experiments, we shall study the polarization in  $B \rightarrow VV$  decays, especially in  $B \rightarrow \rho K^*$  and  $\phi K^*$  decays in this section.

One important point that should be noted is that the predictions for the branching ratios of  $B \rightarrow \rho K^*$  and  $\phi K^*$  modes are well consistent with the experiment data, which means that if we want to solve the observed polarization anomaly, we need to find some way to reduce the longitudinal amplitude and enhance transverse ones simultaneously. Many studies have been made to try to provide possible resolutions to the anomaly both within the SM [36, 37, 38, 39] and in various new physics models [40, 41, 42]. Here we only concentrate on the longitudinal polarization fraction and the main results are listed in Table X.

It is noted that the polarization anomaly could be well resolved by introducing the tensor operators  $O_{T1} = \bar{s}\sigma^{\mu\nu}(1+\gamma^5)b\bar{s}\sigma_{\mu\nu}(1+\gamma^5)s$  and  $O_{T8} = \bar{s}_i\sigma^{\mu\nu}(1+\gamma^5)b_j\bar{s}_j\sigma_{\mu\nu}(1+\gamma^5)s_i$  in Ref. [42]. It is interesting to see that these two operators have similar forms as  $Q_{15}$  and  $Q_{16}$  in Eq. (11). However, from the numerical results given by Table X, we can see that the predicted longitudinal polarization fraction  $f_L$  for these decay modes in the type III 2HDM is almost the same as the one within the SM. Although there are new operator

TABLE VIII: The time-dependent CP asymmetry parameters  $C_f$  (first line) and  $S_f$  (second line) for  $B_d$  decays both in the SM and in the type III 2HDM. Case A-C stand for the three different parameter spaces listed in Section III.

Decay modes	Case A	Case B	Case C	SM
$B^0 \rightarrow \rho^+ \rho^-$	0.043	0.043	0.042	0.035
	-0.95	-0.95	-0.95	-0.95
$B^0 \rightarrow \rho^0 \rho^0$	-0.18	-0.22	-0.12	-0.15
	0.97	0.92	0.99	0.89
$B^0 \rightarrow \omega \rho^0$	0.063	0.063	0.063	0.029
	-0.61	-0.61	-0.62	-0.97
$B^0 \rightarrow \phi \rho^0$	0.011	0.011	0.011	0.014
	0.70	0.70	0.70	0.70
$B^0 \rightarrow \omega \phi$	0.011	0.011	0.011	0.014
	0.70	0.70	0.70	0.70
$B^0 \rightarrow \omega \omega$	0.12	0.09	0.14	0.21
	0.53	0.65	0.40	-0.18
$B^0 \rightarrow K^{*0} \bar{K}^{*0}$	0.092	0.061	0.13	0.15
	0.85	0.92	0.75	0.57

TABLE IX: The time-dependent CP asymmetry parameters  $C_f$  (first line),  $S_f$  (second line), and  $D_f$  (third line) for  $B_s$  decays both in the SM and in the type III 2HDM

Decay modes	Case A	Case B	Case C	SM
$B_s \rightarrow \phi \rho^0$	-0.005	-0.005	-0.005	-0.004
	0.052	0.052	0.052	0.14
	0.99	0.99	0.99	0.99
$B_s \rightarrow \phi \omega$	-0.020	-0.020	-0.020	-0.018
	0.23	0.23	0.23	0.49
	0.97	0.97	0.97	0.87
$B_s \rightarrow \phi \phi$	-0.054	-0.090	-0.060	-0.020
	0.33	0.49	0.14	-0.004
	0.94	0.87	0.99	1.0
$B_s \rightarrow K^{*+} K^{*-}$	-0.30	-0.30	-0.28	-0.20
	0.92	0.95	0.88	0.79
	0.25	0.12	0.39	0.57
$B_s \rightarrow \bar{K}^{*0} K^{*0}$	-0.085	-0.12	-0.049	-0.018
	0.31	0.45	0.15	-0.003
	0.95	0.88	0.99	1.0

contributions in  $B \rightarrow \phi K^*$  modes, we still can not resolve the polarization anomaly observed in this decay mode. This is due to the fact that the strength of new operators in 2HDM is severely suppressed by the factor  $m_q \lambda_{qq}/m_b$ . Moreover, as has already been mentioned in the beginning of this section, the Wilson coefficients of these new operators are very small, which also result in the small effects on the transverse amplitudes.

TABLE X: The longitudinal polarization fractions  $f_L$  for  $B \rightarrow \rho K^*$  and  $\phi K^*$  decay modes. Case A-C stand for the three different parameter spaces in the type III 2HDM.

Decay modes	SM	Case A	Case B	Case C	Exp.
$B^+ \rightarrow \rho^+ K^{*0}$	0.91	0.91	0.91	0.91	$0.48 \pm 0.08$
$B^0 \rightarrow \rho K^{*0}$	0.95	0.95	0.93	0.93	$0.57 \pm 0.12$
$B^+ \rightarrow \phi K^*$	0.89	0.89	0.89	0.89	$0.50 \pm 0.05$
$B^0 \rightarrow \phi K^{*0}$	0.89	0.89	0.89	0.89	$0.491 \pm 0.032$

For the other  $B \rightarrow VV$  decay modes, the predictions for longitudinal polarization fractions are always about  $0.90 \sim 0.95$ . For simplify, we shall not list the results in details anymore.

In conclusion, adopting the current parameter spaces and with the general factorization method, we could not resolve the polarization anomaly observed in  $B \rightarrow \rho K^*$  and  $\phi K^*$  modes within the SM and 2HDM.

## V. CONCLUSIONS

Using the general factorization approach, we have studied all the  $B \rightarrow VV$  decay modes except for pure annihilation decay channels both within the SM and in the two-Higgs-doublet model. From the numerical results given in the previous section, we can see that: for the branching ratios, our predictions are generally well consistent with the current experimental data expect for the  $B_s \rightarrow \phi\phi$  decay mode, and the new physics has margin or even negligible effects on this observable. However, the new physics can give remarkable contributions to the CP asymmetry parameters  $C_f$  and  $S_f$ , especially to  $S_f$  in the penguin-dominated decay modes. Unfortunately, our predictions for the longitudinal polarization fractions of  $B \rightarrow \rho K^*$  and  $\phi K^*$  decay modes in 2HDM are still as large as the ones in the SM, which are much larger than the experimental data. Some new mechanisms may be needed to improve those discrepancies.

For simplicity, in this paper we have neglected the contributions from annihilation and exchange diagrams, although they may play a significant rule in some decay channels. In our numerical calculations, we have only considered three possible parameter spaces for the type III 2HDM. Also we have totally neglected the first generation Yukawa couplings and the off-diagonal matrix elements of the Yukawa coupling matrix, in order to eliminate the FCNC at tree level. However, it is possible that the FCNC involving the third generation quarks still exists at tree level, making the constraints less stronger. In a word, we do not exclude the possibility to improve the predictions by using the other factorization

methods with the annihilation and exchange diagram contributions included, by choosing other parameters spaces, or even by introducing additional fourth-generation quarks [43].

In conclusion, we have shown that the new Higgs bosons in the type III 2HDM with spontaneous CP violation can have significant effects on some charmless  $B \rightarrow VV$  decays, especially for the penguin-dominated decay modes, which can be used as good signals to test the SM and to explore new physics from more precise measurements in the future  $B$ -factory experiments.

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